Features of direct and sequential Coulomb breakup of ^{6}Li ions

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Coulomb dissociation of ^{6}Li in the field of ^{208}Pb at different energies via resonance and continuum levels is discussed in detail. Relations are given which can be used to directly relate the Coulomb breakup cross section to the astrophysical S factor. Predictions for energy dependence and angular distributions are given. The direct Coulomb breakup of ^{6}Li is found to be of the same order of magnitude as the sequential breakup at higher projectile energies. The effect to elastic scattering can be accounted for by introducing a dynamic polarization potential.

I. INTRODUCTION

Breakup processes of nuclear projectiles under the influence of the differential Coulomb field of heavy nuclei are of considerable interest since they provide information on electromagnetically induced interactions of the projectile constituents.[1,2] Experimentally, the situation of pure Coulomb breakup can be realized either by scattering at energies below the Coulomb barrier or, at higher energies for collisions with small deflection angles guaranteeing sufficiently large impact parameters beyond the range of the nuclear interaction. The latter approach has recently been analyzed[3,4] demonstrating interesting possibilities for studies of astrophysical aspects. The breakup may result either from Coulomb transitions to free continuum states of the fragments or from transitions via resonance states above the breakup threshold, followed by a subsequent disintegration into fragments. This (resonant) *sequential breakup* has been found to be dominant at lower projectile energies,[5,6] while the extent to which a "two-step mechanism" contributes at higher energies is not extensively studied. Experimental observations of the $^6\text{Li}\rightarrow\alpha+d$ breakup at projectile energies of 10–30 MeV/amu (Refs. 7–11) seem to indicate that in these cases also, a considerable fraction of the Coulomb breakup cross section has to be attributed to sequential processes via resonant states in ^{6}Li, in particular via the first excited state at E_x ($I^\pi=3^+$)=2.19 MeV. Nevertheless, the investigation of the direct Coulomb breakup mode[12] appears to be interesting. While sequential processes are expected to be well described such as Coulomb excitations of bound states[13] with the lifetime of the resonances larger than the collision time, the direct process involves energy-dependent transition matrix elements into the continuum of the fragment states, distorted by the Coulomb field present at the breakup point. We may expect that for energetic particles the internal distortion of the relative system of the fragments is small. Indeed, with this assumption Coulomb breakup reactions of ^{7}Li have recently been fairly well described[14,15] by using the formalism of Coulomb excitation of quasibound states; i.e., for cases where the actual excitation region and the breakup point are rather distant from each other, the latter far away from the field of the target nucleus.

A reliable and sufficiently accurate theoretical description of experimentally observed direct Coulomb breakup processes would provide an interesting access to those nuclear transition matrix elements which also determine the time-reversed process of breakup, the fusion of the fragments to the projectile nucleus at very low relative energies. In fact, such studies would enable an experimental extension and an important check of the critical ingredients of capture reaction studies at astrophysical energies.[3,4,14]

In context of current experimental investigations[8–10] of Coulomb dissociation of 156 MeV ^{6}Li ions, the present paper explores various features of Coulomb breakup from a theoretical point of view on the basis of a distorted-wave Born approximation (DWBA) approach and a semiclassical approach to the process. The case of ^{6}Li is of particular interest, since studies of the $d(\alpha,\gamma)\,^6$Li capture cross section[16] provide independent information on the electromagnetic transition probabilities for relative energies above 1 MeV, in addition to the $B(E2,1^+\rightarrow 3^+)$ value[17] for the resonance transition at $E_{\alpha d}=0.71$ MeV. Furthermore, the ^{6}Li case is governed by an electromagnetic quadrupole transition, which turns out to be enhanced in the breakup process[18,19] and leads to comparatively large cross sections. We study the specific case of

$$^6\text{Li}+A\rightarrow\alpha+d+A \tag{1.1}$$

in some detail to understand various sensitivities of the cross section.

We evaluate the energy dependence and the angular distribution of the breakup of ^{6}Li scattered off ^{208}Pb via the resonant and continuum states and give relations which relate the breakup cross section with the astrophysical S factor.

II. THEORETICAL CONCEPT

There are various attempts to describe breakup processes.[20–25] The most elaborate theory accounting for the absorption and the distortion by the nuclear field is the post-form DWBA theory worked out by Baur *et al.*[22,23] The theory rests upon a zero-range approximation which constrains the internal momentum distribution of the cluster fragments to a Lorentzian shape with parameter values fixed by the binding energy. Even though the square of the momentum-space wave function of ^{6}Li does not have this shape, the differences between different types of relative motion wave functions are not very strongly developed in the region of small relative momenta $k \leq 0.3$ fm^{-1}. Thus even with a zero-range approximation, the post-form DWBA should be applicable to those cases of the Coulomb breakup where the other important assumption of the post-form DWBA theory, namely, the neglibility of the final-state interaction between the fragments, remains valid. This would happen, obviously, for cases involving high relative energies of the fragments, which, however, are not favored by Coulomb processes. However, cases involving low relative energies are especially favored by Coulomb excitation processes due to the condition of adiabaticity, but for these cases the neglect of the final-state interaction between the fragments cannot be justified.

Thus, at lower relative energies, due to the increased importance of the interaction of fragments in the final state, the (elastic) breakup of the projectile a into the fragments b and x may be considered as a quasisequential process

$$a+A \rightarrow a^* + A \rightarrow b+x+A \tag{2.1}$$

represented by the prevailing prior-form DWBA transition amplitude[21]

$$T_{fi} = \langle X_{Qf}^{(-)}(\mathbf{R})\phi_{\mathbf{k}}^{(-)} \mid U_{bA}+U_{xA}-U_{aA} \mid X_{Qi}^{(+)}(\mathbf{R})\phi_a \rangle \tag{2.2}$$

which is better suited for the studies of the Coulomb dissociation of light ions.[26] It is worthwhile to recall here that the above T matrix *does not* provide a good description of the *nuclear* breakup when the relative energy of the fragments is small, and hence, in the present work we make a special effort to apply it to only those cases where the Coulomb breakup dominates. Here

$$X_{\mathbf{Q}i}^{(+)}(\mathbf{R})$$

and

$$X_{\mathbf{Q}f}^{(-)}(\mathbf{R})$$

denote the center-of-mass motion of the initial and the final state with the momenta $\mathbf{Q}_i$ and $\mathbf{Q}_f$, respectively. The wave functions $\phi_a(\mathbf{r})$ and $\phi_{\mathbf{k}}(-)(\mathbf{r})$ represent the ground state and the continuum ("excited") state of the relative motion of the projectile. When the fragments b and x are observed with the momenta $\mathbf{k}_b$ and $\mathbf{k}_x$, the momenta are given by

$$\mathbf{Q}_f = \mathbf{k}_b + \mathbf{k}_x \ , \tag{2.3a}$$

$$\mathbf{k} = \frac{m_b}{m_a}\mathbf{k}_x - \frac{m_x}{m_a}\mathbf{k}_b \ . \tag{2.3b}$$

In the case of a pure sequential process $\phi_{\mathbf{k}}$ is a resonance state having a substantial overlap with ground-state wave function, but the same matrix element is also expected to describe nonresonant breakup processes when adequate wave functions for the continuum are introduced.

Assuming point-charge distributions for the constituent projectile clusters, the residual Coulomb interaction for $R > r$ is

$$V_{\text{res}} = Z_T e^2 \left[\frac{Z_b}{r_b} + \frac{Z_x}{r_x} - \frac{Z_a}{R} \right] \tag{2.4}$$

$$= 4\pi Z_T e^2 \sum_{L,M\geq 1} \left[Z_b \left[-\frac{m_x}{m_a} \right]^L + Z_x \left[\frac{m_b}{m_a} \right]^L \right] \times \frac{r^L}{R^{L+1}} \frac{1}{2L+1} Y_{LM}^*(\mathbf{R}) Y_{LM}(\mathbf{r}) \ . \tag{2.5}$$

The cross section for the breakup of the projectile, when the center-of-mass wave vector of the fragments lies between $\mathbf{Q}_f$ and $\mathbf{Q}_f + d\mathbf{Q}_f$, the relative motion wave vector lies between $\mathbf{k}$ and $\mathbf{k}+d\mathbf{k}$, and the spin orientations are unspecified, is given by[27]

$$d\sigma = \frac{2\pi}{\hbar v}(2I_i+1)^{-1} \sum_{M_i M_f} |T|^2 \frac{d\mathbf{Q}_f d\mathbf{k}}{(2\pi)^6} \delta(E_i - E_f) \ , \tag{2.6}$$

which yields for Lth multipole

$$\frac{d^3\sigma_{EL}}{d\Omega_{Q_f} d\Omega_k d\varepsilon} = \frac{4(Z_T e)^2}{\hbar v (2L+1)^2} \frac{\mu Q_f}{\hbar^2} \cdot \frac{\mu_{bx} k}{(2\pi)^3 \hbar^2} \cdot \left[Z_b e \left[-\frac{m_x}{m_a} \right]^L + Z_x e \left[\frac{m_b}{m_a} \right]^L \right]^2 \times (2I_i+1)^{-1} \sum_{M_i M_f} \cdot \left| \sum_M \langle \phi_{\mathbf{k}}^{(-)}(\mathbf{r}) \mid r^L Y_{LM}(\mathbf{r}) \mid \phi_a(\mathbf{r}) \rangle \langle X_{Q_f}^{(-)}(\mathbf{R}) \mid R^{-L-1} \cdot Y_{LM}^*(\mathbf{R}) \mid X_{Q_i}^{(+)}(\mathbf{R}) \rangle \right|^2 , \tag{2.7}$$

where v is the velocity of the projectile in the center-of-mass system. The fragments have a relative energy between ε and $\varepsilon + d\varepsilon$ and the rest of the symbols have their usual meaning. Performing the $d\Omega_k$ integration and using the angular momentum algebra[28] we get

$$\frac{d^2\sigma_{EL}}{d\Omega_{Q_f}\cdot d\varepsilon}=\left[\frac{Z_T e}{\hbar v}\right]^2\cdot\frac{4Q_iQ_f}{(2L+1)^3}\cdot B(EL,\varepsilon)\sum_M\left|\langle X_{Q_f}^{(-)}(\mathbf{R})\,|\,R^{-L-1}Y_{LM}(\mathbf{R})\,|\,X_{Q_i}^{(+)}(\mathbf{R})\rangle\right|^2 . \tag{2.8}$$

This can be rewritten to give the Coulomb dissociation cross section

$$\frac{d^2\sigma_{EL}}{d\Omega\, d\varepsilon}=\left[\frac{Z_T e}{\hbar v}\right]^2\cdot d^{-2L+2}\cdot B(EL,\varepsilon)\frac{df_{EL}}{d\Omega} , \tag{2.9}$$

where we have dropped the subscript Q_f from the solid angle into which the center of mass of the fragmented projectile is scattered. Additionally, we have defined d as half the distance of minimum approach,

$$d=\frac{Z_a\cdot Z_T\cdot e^2}{2E} \tag{2.10}$$

and $B(EL,\varepsilon)$ is the reduced transition probability per unit energy for the transition from the bound (ground) state of the projectile to the continuum state for the transition having the multipolarity 2^L,

$$B(EL,\varepsilon)=\frac{\mu_{bx}k}{(2\pi)^3\hbar^2}\left[Z_b\left[-\frac{m_x}{m_a}\right]^L\cdot e+Z_x\left[\frac{m_b}{m_a}\right]^L e\right]^2 \times\sum_{M_f M}\int|\langle\phi_{\mathbf{k}}(\mathbf{r})\,|\,r^L Y_{LM}(\mathbf{r})\,|\,\phi_a(\mathbf{r})\rangle|^2 d\Omega_k \tag{2.11}$$

where the factor $[\mu_{bx}k/(2\pi)^3\cdot\hbar^2]$ ensures the energy normalization of the final-state wave function $\phi_{\mathbf{k}}$. The initial and the final states have spins (I_i,M_i) and $(\mathbf{I}_f,\mathbf{M}_f)$, respectively. The Coulomb excitation function

$$\frac{df_{EL}}{d\Omega} ,$$

is given by[28]

$$\frac{df_{EL}}{d\Omega}=\frac{4Q_iQ_f}{(2L+1)^3}\cdot d^{2L-2} \times\sum_M\left|\langle X_{Q_f}^{(-)}(\mathbf{R})\,|\,R^{-L-1}Y_{LM}(\mathbf{R})\,|\,X_{Q_i}^{(+)}(\mathbf{R})\rangle\right|^2 . \tag{2.12}$$

The reduced transition probability per unit energy $B(EL,I_i\rightarrow I_f,\varepsilon)$ relevant for our case is related to the transition probability $B_{\text{cap}}(EL,I_f\rightarrow I_i,\varepsilon)$ for capture from the state $|\mathbf{I}_f\mathbf{M}_f\rangle$ with relative energy ε, to the ground state $|I_iM_i\rangle$ by

$$B_{\text{cap}}(EL,I_f\rightarrow I_i,\varepsilon)=\frac{\pi^2\hbar^3}{\mu_{bx}\varepsilon}\frac{(2I_i+1)}{(2I_f+1)}B(EL,I_i\rightarrow I_f,\varepsilon) , \tag{2.13}$$

where the capture cross section is given by[29]

$$\sigma_{\text{cap}}(EL,I_f\rightarrow I_i,\varepsilon) =\frac{8\pi(L+1)}{L(2L+1)!!^2}\cdot\frac{1}{\hbar}\left[\frac{E_Y}{\hbar c}\right]^{2L+1}B_{\text{cap}}(EL,I_f\rightarrow I_i,\varepsilon) , \tag{2.14}$$

where $E_\gamma(=\hbar \mathbf{k}_\gamma c)$ is the energy of the gamma ray emitted after the capture. Noting that the capture cross section for low relative energies is related to the astrophysical S factor by

$$S(\varepsilon)=\varepsilon\cdot\sigma_{\text{cap}}(\varepsilon)\cdot e^{2\pi\eta},\quad \eta=\frac{Z_b\cdot Z_x e^2}{\hbar v_{bx}} , \tag{2.15}$$

we may express

$$B(EL,I_i\rightarrow I_f,\varepsilon)=\frac{L(2L+1)!!^2}{8\pi(L+1)}\frac{(2I_f+1)}{(2I_i+1)} \times\frac{\mu_{bx}}{\pi^2\hbar^2}\cdot\frac{S(\varepsilon)\cdot e^{-2\pi\eta}}{k_\gamma^{2L+1}} \tag{2.16}$$

in units of fm^{2L+1}, which can be converted for normal units (of $e^2\text{fm}^{2L}/\text{MeV}$) by dividing it by e^2.

The reduced transition probability in the case of a narrow resonance is obtained by the integration

$$B(EL,I_i\rightarrow I_f)=\int_{\text{Res}}B(EL,\varepsilon)d\varepsilon , \tag{2.17}$$

while the reduced transition probability for the continuum transition with unspecified spins is analogously defined as

$$B(EL,\text{cont})=\sum_{I_f}\int_{\text{cont}}B(EL,\varepsilon)d\varepsilon . \tag{2.18}$$

In order to obtain a reliable estimate of the angle integrated Coulomb dissociation cross section, we realize that such trajectories which lead to distances less than the sum of the two nuclear radii, will be strongly absorbed, and those which stay clear of this distance will feel only the Coulomb interaction. Thus semiclassically, one can obtain the pure Coulomb dissociation cross section of the projectile by integrating the differential cross section up to the angle θ_c (Refs. 28 and 30) where

$$\sin\frac{\theta_c}{2}=\frac{1}{2E/E_B-1} \tag{2.19}$$

and

$$E_B=Z_a\cdot Z_T\cdot e^2/R_{\text{cut}} \tag{2.20}$$

with

$$R_{\text{cut}}=1.36(A_a^{1/3}+A_T^{1/3})+0.5 . \tag{2.21}$$

For comparison the reaction cross section σ_R for nucleus-nucleus collisions can be estimated as

$$\sigma_R = \pi R_{\rm cut}^2 \left[1 - \frac{E_B}{E}\right], \tag{2.22}$$

which has been found to give a good description of the experimental data with the choice (2.21) of $R_{\rm cut}$.[31] These semiclassical relations can be translated to the quantum-mechanical description by introducing a lower cutoff[26]

$$L_c = Q_i R_{\rm cut} \left[1 - \frac{2\eta_i}{Q_i R_{\rm cut}}\right]^{1/2} \tag{2.23}$$

of the orbital angular momentum in a partial-wave expansion for $df_{EL}/d\Omega$ above.

We have found that for the cases to be reported later, the semiclassical and the quantal descriptions of the Coulomb-excitation function $df_{EL}/d\Omega$ agree to better than 0.5%. The energy-differential Coulomb-dissociation cross section is obtained

$$\frac{d\sigma_{EL}}{d\varepsilon} = \left[\frac{Z_T e}{\hbar v}\right]^2 d^{-2L+2} B(EL,\varepsilon) f_{EL}^N, \tag{2.24}$$

where the superscript N over the total Coulomb-excitation function f_{EL}^N denotes the nuclear absorption with

$$f_{EL}^N = \int_{\theta=0}^{\theta=\theta_c} \frac{df_{EL}}{d\Omega} \cdot d\Omega \tag{2.25}$$

semiclassically, and

$$f_{EL}^N = \frac{64\pi^2}{(2L+1)^3} \cdot Q_i Q_f d^{2L-2} \times \sum_{L_i L_f > L_c} (2L_i+1)(2L_f+1) \times \begin{bmatrix} L_i & L_f & L \\ 0 & 0 & 0 \end{bmatrix} | M_{L_i L_f}^{-L-1} |^2 \tag{2.26}$$

in the corresponding quantum-mechanical description.

In the above

$$M_{L_i L_f}^{-L-1}$$

are the well-known Coulomb-excitation matrix elements which we evaluate in the well fulfilled WKB approximation[28] due to the lower cutoff L_c, which limits the angular momenta to large values.

Now the Coulomb-dissociation cross section is obtained as

$$\sigma_{EL} = \left[\frac{Z_T e}{\hbar v}\right]^2 \cdot d^{-2L+2} \cdot \int B(EL,\varepsilon) f_{EL}^N d\varepsilon . \tag{2.27}$$

As f_{EL} is a slowly varying function of the adiabaticity parameter,

$$\xi = n_f - n_i , \tag{2.28}$$

we can take it out of the integration sign in the above for *narrow resonances,* thus writing

$$\sigma_{EL}^{\rm Res}(I_i \to I_f) = \left[\frac{Z_T e}{\hbar v}\right]^2 \cdot d^{-2L+2} \cdot B(EL, I_i \to I_f) f_{EL}^N(\xi_{\rm res}) , \tag{2.29}$$

where the experimentally known values of $B(EL)$ (in units of $e^2{\rm fm}^{2L}$) can be used.

In order to estimate the relative importance of the Coulomb-dissociation processes, the nuclear breakup cross section is evaluated by use of the simple and yet reliable Serber model[20] to get the nonelastic breakup (also called absorptive stripping) as

$$\sigma_{\rm as} = \frac{\pi}{2} \cdot R_T \cdot R_P , \tag{2.30}$$

where we have taken R_T as equal to $R_{\rm cut}$ and

$$R_P = \frac{1}{2a}, \quad a = \left[\frac{2\mu_{bx} BE}{\hbar^2}\right]^{1/2} . \tag{2.31}$$

The elastic breakup, due to the nuclear field, which will automatically be distinguished from the nonelastic breakup in a coincidence measurement is akin to the "free dissociation" discussed by Glauber[32] and its cross section is given by

$$\sigma_{\rm fd} = \frac{\pi}{3}(2\ln 2 - \tfrac{1}{2}) R_T \cdot R_P \tag{2.32}$$

$$\simeq 0.59 \sigma_{\rm as} . \tag{2.33}$$

The total breakup cross section is then a sum of $\sigma_{\rm as}, \sigma_{\rm fd}$, and the Coulomb breakup cross section.

III. CLASSICAL TRAJECTORY CONSIDERATIONS

In order to isolate the region where the Coulomb processes may dominate we have performed classical trajectory calculations for the scattering of the ^{6}Li$+^{208}$Pb system at 30, 60, 90, and 156 MeV incident energies, where the projectile moves under the influence of Coulomb and nuclear potential, the latter given by[33]

$$V_N(R) = -240 \cdot \left[1 + \exp\left[\frac{R - 1.17 A_T^{1/3}}{0.766}\right]\right]^{-1} . \tag{3.1}$$

We give a plot of the deflection functions for the above cases in Fig. 1. At 30 MeV, the incident ^{6}Li ion does not see the nucleus at all and its motion is governed completely by the Coulomb field. The other cases describe the situation above the Coulomb barrier. The most remarkable feature which emerges from this is the disappearance of the well-pronounced hump in the deflection function near the Coulomb rainbow angle and onset of orbiting at an incident angular momentum, very close to the rainbow "partial wave." This has an important and beneficial consequence for our study indicating a possibility to isolate an angular range in breakup experiments as a "nuclear free zone." In Table I we have compiled the angular momentum values leading to certain angles of scattering along with the corresponding distances of minimum approach. We see, e.g., that for a scattering of the center of mass of the projectile to 9° the strong ab-

TABLE I. Nature of contributions for trajectories leading to a specific scattering angle in $^{6}Li+^{208}Pb$ collisions at different energies, in a purely classical calculation.

Scattering angle deg	E MeV	l values $\hbar$	d_{min}[a] fm	Remarks
9.0	30	219	83.5	Pure Coulomb
	60	2	0.3	Absorbed
		26	3.2	Absorbed
		150	39.8	Pure Coulomb
	90	4	0.6	Absorbed
		36	4.2	Absorbed
		130	28.6	Pure Coulomb
	156	12	1.3	Absorbed
		66	7.0	Absorbed
		78	13.5	Weakly nuclear
		94	15.9	Pure Coulomb
5.0	156	7	0.7	Absorbed
		66	7.0	Absorbed
		71	12.0	Absorbed
		172	27.5	Pure Coulomb
3.0	156	4	0.4	Absorbed
		66	7.0	Absorbed
		289	45.8	Pure Coulomb

[a]d_{min} is the distance of minimum approach.

sorption by the nucleus would lead to a relatively pure Coulomb contribution at 30, 60, and 90 MeV, whereas the data at 156 MeV will not be free from the nuclear contribution. The data more forward to 5° at 156 MeV are seen to be free from nuclear contribution. The R_{cut} defined by (2.21) above gives a lower limit of the distance of minimum approach beyond which the scattering process may be induced only by the Coulomb field.

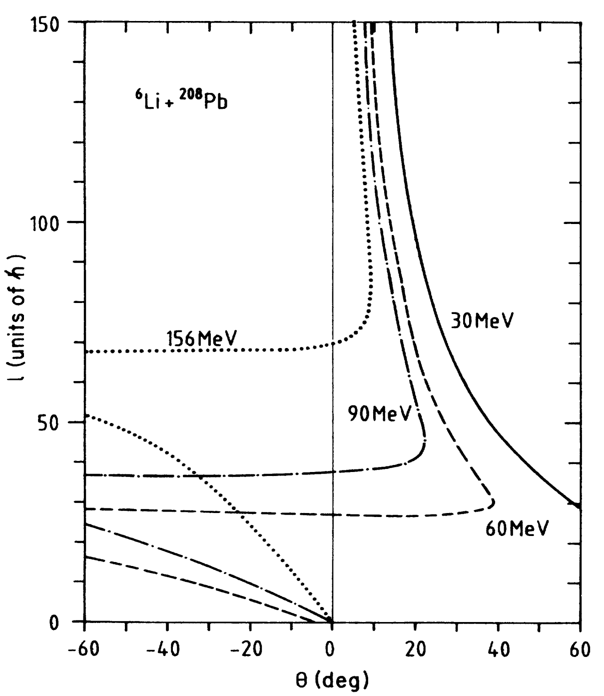

FIG. 1. Classical deflection function for $^{6}Li+^{208}Pb$ system under the influence of nuclear and Coulomb potential at 30, 60, 90, and 156 MeV.

IV. RESULTS OF MODEL CALCULATIONS

The quantity of fundamental importance in these studies is the reduced transition probability per unit energy, which is directly related to the fusion cross section data [Eqs. (2.13)–(2.16)]. It is quite clear that the $L=2$ transitions will dominate the Coulomb dissociation of $^{6}Li\rightarrow\alpha+d$, and in the following we give results only for this case. ^{6}Li has three well-developed resonances having $L=2$, and $I^{\pi}=3^{+}$, 2^{+}, and 1^{+} above the particle emission threshold (see Table II).

In Fig. 2 we have plotted the reduced transition probability per unit energy as a function of the relative energy (ε) of the $\alpha+d$ system for the above-mentioned resonances and the continuum. For the resonances states we

TABLE II. Reduced electromagnetic transition probabilities for resonant and continuum levels of ^{6}Li, having excitation energies $E_x=\varepsilon+1.47$.

I_f	ε (MeV)	Γ (MeV)	$B(E2)$ (e^2fm^4)	Ref.
3^+	0.71	≈ 0.02	24.	a
2^+	3.05	≈ 0.60	17.14	b
1^+	4.03	≈ 1.00	10.29	b
Σ_{I_f}	Continuum		51.66	c

[a]Reference 18.
[b]Obtained by using $B(E2,I_i\rightarrow I_f')=(2I_f'+1)B(E2,I_i\rightarrow I_f)/(2I_f+1)$, with $I_f=3$ (following Ref. 27).
[c]By using Eq. (2.18) including the spectroscopic factor $C^2S\approx 0.5$ for $^{6}Li\rightarrow\alpha+d$ configuration in the $2S$ ground state.

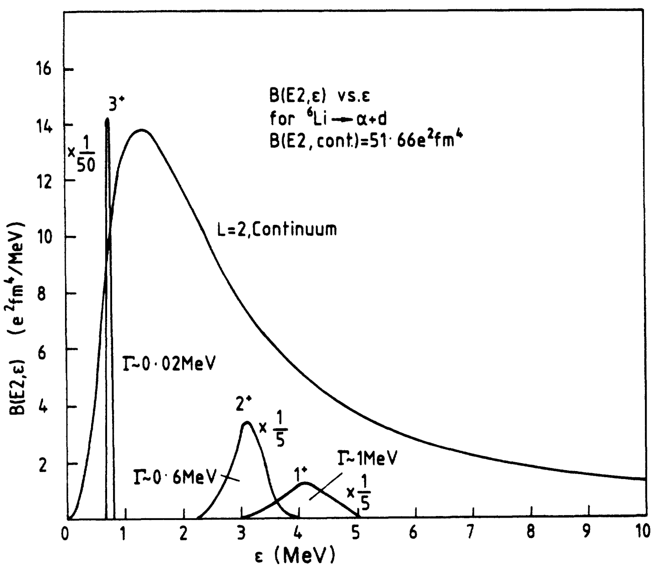

FIG. 2. Reduced transition probability per unit energy for transitions to continuum levels in the resonance and the nonresonant regions for $^6\mathrm{Li}\rightarrow\alpha+d$ breakup.

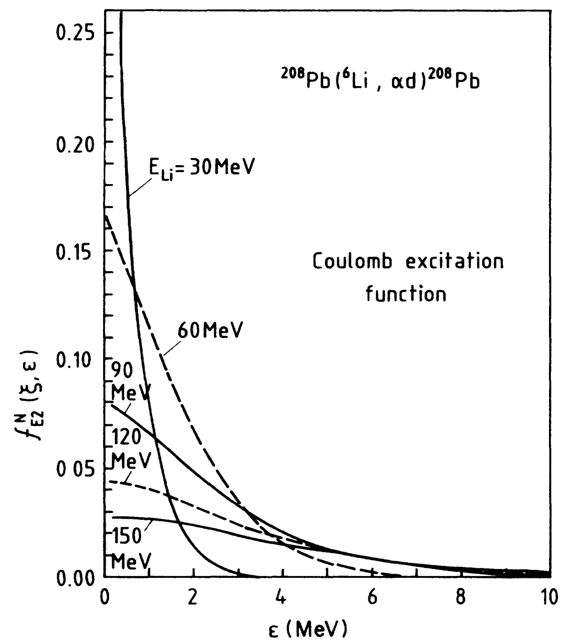

FIG. 3. Coulomb-dissociation (-excitation) function with nuclear absorption for $^6\mathrm{Li}+^{208}\mathrm{Pb}$ system at 30, 60, 90, 120, and 150 MeV as a function of relative energy (ε) of the fragments.

have taken

$$B(E2,\varepsilon)=\left|\frac{i\Gamma/2}{(\varepsilon-\varepsilon_{\mathrm{res}})+i\Gamma/2}\right|^2\cdot B(E2,\varepsilon_{\mathrm{res}}) \qquad (4.1)$$

so that

$$B(E2,I_i\rightarrow I_f,\mathrm{res})=\pi\cdot\frac{\Gamma}{2}\cdot B(E2,\varepsilon_{\mathrm{res}})\;. \qquad (4.2)$$

The corresponding quantity for the continuum is evaluated by calculating ϕ_a and ϕ_k in the nuclear potential

$$V_{ad}(r)=-85.54\cdot\left[1+\exp\left[\frac{r-1.779}{0.70}\right]\right]^{-1}, \qquad (4.3)$$

which reproduces the binding energy of ^{6}Li and its rms radius. The spin-orbit interaction has been neglected for simplicity. This potential is similar to the α-d interaction potential used by Robertson *et al.*[17] for evaluating the S factor for α-d fusion in good agreement with data. A good agreement with breakup experimental results would confirm our belief that such experiments can be used[10] to get reliable experimental information about reaction cross sections of astrophysical interest. It should be remarked here that the rapid increase of $B(E2,\varepsilon)$ from zero as ε increases comes from the small binding energy of ^{6}Li. In addition, as the larger values of ε are strongly suppressed in these processes due to the condition of adiabaticity, this causes a large Coulomb dissociation. At very high incident energies the condition of adiabaticity is quite relaxed and $B(E2,\mathrm{cont})$ defined by (2.18) becomes a measure of breakup via continuum.

The condition of adiabaticity referred to above can be understood by inspecting Fig. 3 where the total Coulomb-excitation function f_{E2}^N is plotted as a function of the breakup energy ε. At low incident energies only the smaller values of the breakup energies ε are important whereas, with increasing incident energy, states with larger values of ε are equally well excited.

In Fig. 4 the energy-differential cross section (2.24) for the Coulomb dissociation of ^{6}Li is shown for various incident energies. It is interesting to note that the relative

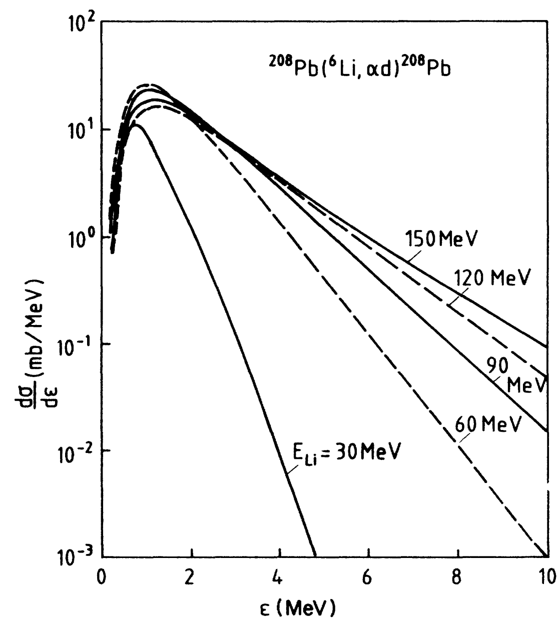

FIG. 4. Energy differential cross section for direct Coulomb dissociation of ^{6}Li scattered off ^{208}Pb at 30, 60, 90, 120, and 150 MeV.

energy spectra of direct Coulomb breakup show a broad peak which tends to get asymmetric and broadened with increasing incident energy. This peak misconstrued as a "resonance" in Ref. 34 is just the result of the superposition of the energy dependence of the $B(EL,\varepsilon)$ values (dominated by the Coulomb penetration) and of the Coulomb-excitation function f_{EL}^{N}, which reflects the equivalent photon spectrum inducing the breakup (see Refs. 3 and 9). This feature is also seen in Fig. 5 where the double-differential cross section for the 3^+ resonance and for the continuum levels integrated up to $\varepsilon=5$ MeV are plotted when the center of mass of the fragments is scattered from ^{208}Pb to $\theta=3°$.

In order to get a feeling of the range of the partial wave contributing to the Coulomb dissociation of ^{6}Li, at different energies, in Fig. 6 we have plotted the partial Coulomb-dissociation cross section (σ_l) as a function of the incident angular momentum (l) for breakup via the 3^+ resonance at 30, 90, and 156 MeV [see Eqs. (2.26) and (2.29)]. The dashed curves give the variation of σ_l when nuclear absorption is absent. The solid curves describe the situation when the nuclear absorption is accounted for [Eq. (2.26)]. We see that the Coulomb-dissociation probability is a slowly varying function of the incident angular momentum and that a very large number of partial waves contribute to the Coulomb breakup of light ions. This aspect enhances the value of the WKB approximation utilized in the present work for evaluation of the radial matrix elements.

Figure 7 displays the differential cross section for the Coulomb breakup of ^{6}Li scattered off ^{208}Pb at 156 MeV

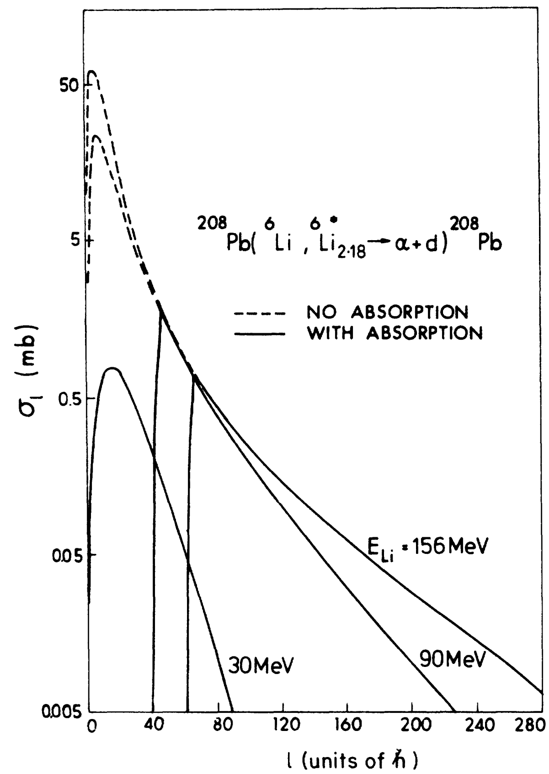

FIG. 6. Partial Coulomb-dissociation cross section via the first 3^+ state of ^{6}Li colliding with ^{208}Pb at 30, 90, and 156 MeV with (——) and without (— — —) nuclear absorption.

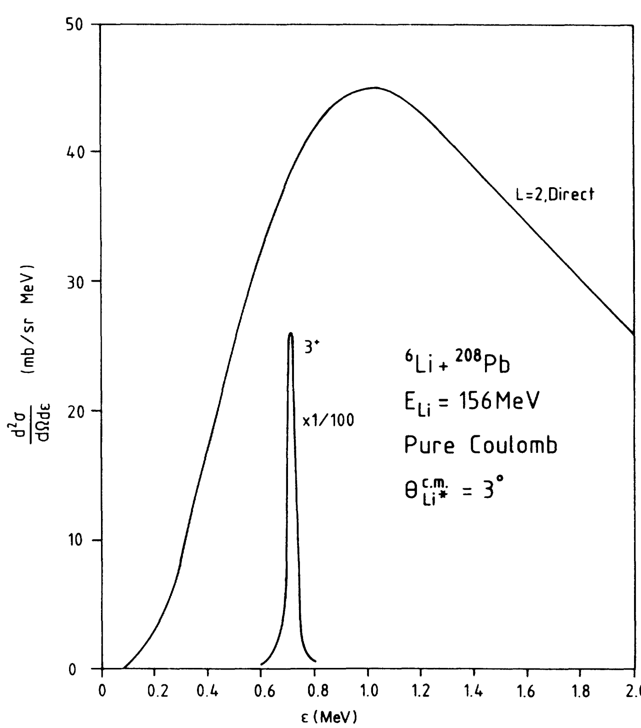

FIG. 5. Double differential cross section for the Coulomb dissociation of ^{6}Li at $\theta_{c.m.}=3°$ via resonant and continuum levels.

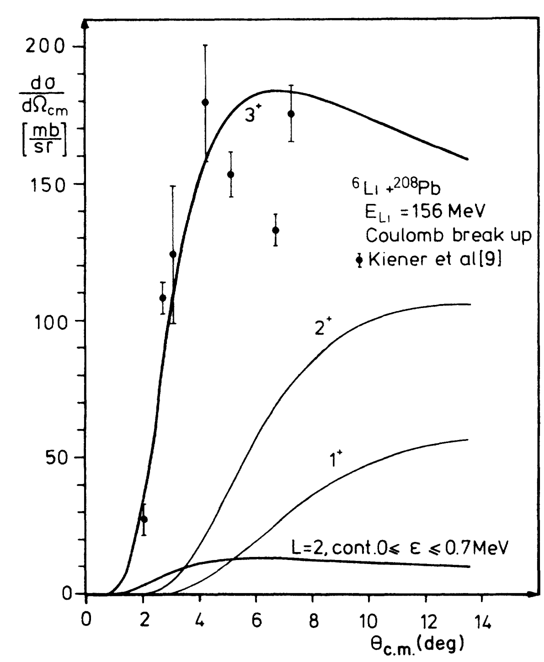

FIG. 7. Predictions for the differential cross section for Coulomb disintegration of ^{6}Li scattered off ^{208}Pb at 156 MeV via 3^+, 2^+, and 1^+ resonant states and via the continuum transitions having $L=2$ and to excitation energies ≤ 0.7 MeV. The experimental data are from Kiener *et al.* (Ref. 9) for breakup via the 3^+ state.

via the 3^+, 2^+, and the 1^+ resonances and the $L=2$ continuum integrated below the 3^+ resonance. For $\theta \leq 5°$ the theoretical prediction is in quantitative agreement with recent experimental results[9] for the case of 3^+ resonance. This is a strong experimental support of our conclusion (see Sec. III) that elastic breakup processes observed at very forward reaction angles and isolated by coincidence measurements from nonelastic events originate from the Coulomb breakup.

Figure 8 shows the energy dependence of the Coulomb dissociation cross section for ^{6}Li scattered off ^{208}Pb for the cases considered in this work. The continuum breakup cross sections are integrated over 0 to 15 MeV and over the continuum energy states below the first resonance. The two experimental data points for $E_{\rm Li}=23$ MeV for the direct and the sequential breakup (via the 3^+ state) are results from Scholz *et al.*[12] The quantitative agreement for the direct breakup without any adjustment of parameters is especially noteworthy. The Coulomb dissociation cross section increases with energy as long as the nuclear absorption can be ignored, which affects more and more partial waves at higher energies. This is evident at energies $E_{\rm Li} \geq 60$ MeV. For comparison Table III gives the cross sections of various types of breakup processes for the particular example of 156 MeV ^{6}Li ions incident on ^{208}Pb. The direct and the sequential Coulomb breakup of ^{6}Li due to the Coulomb processes taken together prove to be about 20% of nuclear breakup. However the direct and the sequential breakup of ^{6}Li due to the Coulomb field are of the same magnitude. At the same time the Coulomb dissociation of ^{6}Li is not a large part of the nuclear reaction cross section. This feature

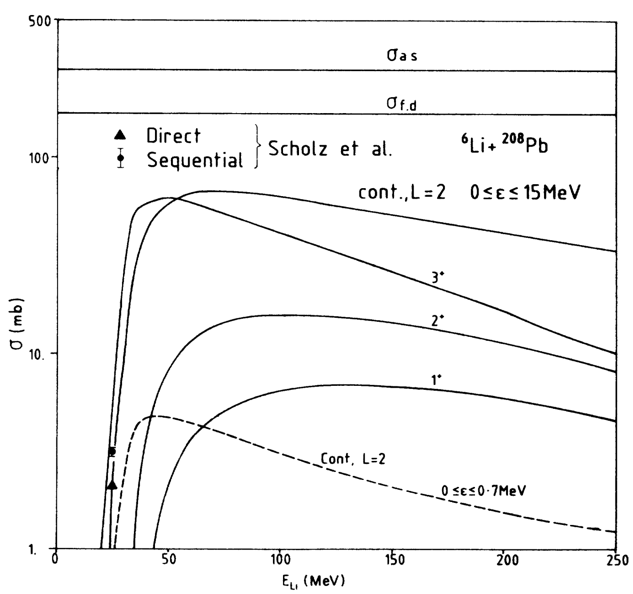

FIG. 8. Energy dependence of the Coulomb dissociation of ^{6}Li scattered off ^{208}Pb via the resonant and continuum states having $L=2$. Nuclear breakup via absorptive stripping ($\sigma_{\rm as}$) and free-dissociation ($\sigma_{\rm fd}$) are also given.

TABLE III. Cross sections for various breakup processes for 156 MeV ^{6}Li($\rightarrow \alpha + d$) incident on ^{208}Pb.

Process of breakup	Model	Cross section
Absorptive stripping	Serber[a]	280 mb
Free-dissociation	Glauber[b]	170 mb
Resonant Coulomb	Present work	
3^+		26 mb
2^+		19 mb
1^+		6 mb
Direct Coulomb	Present work	
$L=2$	$0 \leq \varepsilon \leq 0.7$ MeV	2 mb
	$0 \leq \varepsilon \leq 15$ MeV	50 mb
Reaction cross section		$\simeq 3b$

[a]Reference 20.
[b]Reference 32.

appears to be different from the case of deuteron breakup[35] where the Coulomb dissociation cross section represents a large fraction of the total breakup process as well as of the total reaction cross section. A perturbative approach like DWBA may be less justified in such a case.

Nevertheless, due to the relatively large $B(E2)$ value for the 3^+ resonance and the low dissociation threshold, Coulomb excitation and breakup leads to an additional absorption for elastic scattering of ^{6}Li projectiles, accounted by a long-range imaginary potential. Neglecting the adiabaticity corrections (see Ref. 35) it is given for $R > R_{\rm cut}$ by

$$\mathrm{Im}(U_{pL}, L=2, 1^+ \rightarrow 3^+) = -\frac{\pi^2 Q_i}{4E} \cdot \frac{(Z_T e)^2}{2L+1} \cdot \frac{B(E2, 1^+ \rightarrow 3^+)}{R^{2L+1}} . \qquad (4.4)$$

At higher energies the condition for adiabaticity is quite relaxed and the dynamic polarization potential arising from the Coulomb dissociation of ^{6}Li via $L=2$ states would be obtained by replacing the $B(E2)$ value above by a sum of the corresponding values for the resonant and the continuum states (see Table II).

V. CONCLUDING REMARKS

We have studied the features and the sensitivities of the Coulomb dissociation cross section of ^{6}Li scattered off ^{208}Pb at various energies via the resonant and the continuum levels. The Coulomb dissociation probability is found to be large over a large range of impact parameters and it is not negligible even at the highest energy considered. Whereas the resonant breakup dominates at lower incident energies, the direct breakup is of the same order as the sequential breakup at higher energies. This may lead to an absorptive dynamic polarization potential varying as $1/R^5$ contributing to the imaginary part of the optical-model potential. We feel that the results obtained in the present study are interesting as they can also relate the breakup cross section with fusion cross sections of astrophysical interest.

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